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Compressional properties of nuclear matter in the relativistic mean field theory with the excluded volume effects

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ABSTRACT

Compressional properties of nuclear matter are studied by using the mean field theory with the excluded volume effects of the nucleons. It is found that the excluded volume effects make it possible to fit the empirical data of the Coulomb coefficient K_c of nucleus incompressibility, even if the volume coefficient K is small($\sim$ 150MeV). However, the symmetry properties favor $K = 300 \pm 50$MeV as in the cases of the mean field theory of point-like nucleons.

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One way to determine the incompressibility K of nuclear matter from the giant monopole resonance (GMR) data is using the leptodermous expansion[1] of nucleus incompressibility $K(A,Z)$ as follows.

$$K(A,Z) = K + K_{sf}A^{-1/3} + K_{vs}I^2 + K_c Z^2 A^{-4/3} + \cdots \quad ; \quad I = 1 - 2Z/A, \tag{1}$$

where the coefficients K_{sf}, K_{vs} and K_c are surface term coefficient, volume-symmetry coefficient and Coulomb coefficient, respectively. We have omitted higher terms in eq. (1). Although there is uncertainty in the determination of these coefficients by using the present data, Pearson [2] pointed out that there is a strong correlation among K, K_c and the skewness coefficient, i.e., the third-order derivative of nuclear saturation curve. (See table 1.) Similar observations are done by Shlomo and Youngblood [3].

Table 1

According to this context, Rudaz et al. [4] studied the relation between incompressibility and the skewness coefficient by using the generalized version of the relativistic Hartree approximation [5]. The compressional and the surface properties are studied by Von-Eiff et al. [6][7][8] in the framework of the mean field approximation of the σ-ω-ρ model with the nonlinear σ terms. They found that low incompressibility ($K \approx 200$MeV) and a large effective nucleon mass M_0^* at the normal density ($0.70 \leq M_0^*/M \leq 0.75$) are favorable for the nuclear surface properties [8]. On the other hand, using the same model, Bodmer and Price [9] found that the experimental spin-orbit splitting in light nuclei supports $M_0^* \approx 0.60M$. The result of the generator coordinate calculations for breathing-mode GMR by Stoitsov, Ring and Sharma [10] suggests $K \approx 300$MeV.

In previous papers[11][12], we have studied the effective nucleon mass M_0^*, incompressibility K and the skewness K' in detail, using the relativistic mean field theory with the nonlinear σ terms [13] and the one with the nonlinear σ and ω terms [14]. We found that $K = 300 \pm 50$MeV is favorable to account for K, K_c, K_{vs} and the symmetry energy a_4, simultaneously [11][12]. It was also found that the empirical values of K and K_c in table 1 are not well reproduced by these model, if $K \lesssim 200$MeV [11][12].

On the other hand, the incompressibility K, which is calculated in the framework of the quark-meson coupling (QMC) model is rather small($\sim$ 200MeV) [15][16][17]. The sub-structure or the finite size effect of nucleons may be important in calculating K. In this paper, we study M_0^*, K, K', K_c and K_{vs}, which can be calculated in the framework of nuclear matter with aid of the scaling model [1], by using the relativistic mean field theory with the excluded volume effects (EVE) of nucleons [18][19][20], and compare the results with the GMR data. We also examine whether the QMC result is reproduced by the EVE model.

We use the relativistic mean field theory with the EVE of nucleons [18][19][20]. As a Lagrangian, we use the σ-ω model with the nonlinear σ terms as in ref. [20].

(For a while, we restrict our discussions to the symmetric nuclear matter and do not consider the ρ meson effects.) The Lagrangian density consists of three fields, the nucleon ψ, the scalar σ-meson ϕ, and the vector ω-meson V_μ, i.e.,

$$L_{N\sigma\omega} = \bar{\psi}(i\gamma_\mu\partial^\mu - M)\psi + \frac{1}{2}\partial_\mu\phi\partial^\mu\phi - \frac{1}{4}F_{\mu\nu}F^{\mu\nu} + \frac{1}{2}m_v^2 V_\mu V^\mu$$
$$+g_s\bar{\psi}\psi\phi - g_v\bar{\psi}\gamma_\mu\psi V^\mu - U(\phi) \quad ; \qquad F_{\mu\nu} = \partial_\mu V_\nu - \partial_\nu V_\mu, \tag{2}$$

where m_v, g_s and g_v are ω-meson mass, σ-nucleon coupling and ω-nucleon coupling, respectively. The potential $U(\phi)$ includes a nonlinear cubic-quartic terms of the scalar field ϕ; i.e.,

$$U(\phi) = \frac{1}{2}m_s^2\phi^2 + \frac{1}{3}b\phi^3 + \frac{1}{4}c\phi^4, \tag{3}$$

where m_s is σ-meson mass, and b and c are the constant parameters which are determined phenomenologically. The Lagrangian density (2) is also the same as the one used in the theory of Boguta and Bodmer [13].

In the mean field theory of the point-like nucleon, at zero-temperature, the baryon density ρ_{pt} is given as

$$\rho_{pt} = \frac{\lambda}{3\pi^2}k_F^3, \tag{4}$$

where k_F is Fermi momentum and $\lambda = 2$ in the nuclear matter. In the model with the EVE [18][19][20], the volume V for the N body system of nucleons in configurational space is reduced to the effective one, $V - NV_n$, where V_n is the volume of a nucleon. According to this modification for the volume, the baryon density ρ is given by

$$\rho = \frac{\rho'}{1 + V_n\rho'}, \tag{5}$$

where ρ' has the same expression as ρ_{pt} for the given k_F. In a similar way, the scalar density is given by

$$\rho_s = \frac{\rho_s'}{1 + V_n\rho'} \tag{6}$$

where ρ_s' has the same expression as the scalar density of the system of the point-like nucleons and is given by

$$\rho_s' = \frac{\lambda}{2\pi^2}M^*[k_F\sqrt{k_F^2 + M^{*2}} - M^{*2}\ln(\frac{k_F + \sqrt{k_F^2 + M^{*2}}}{M^*})], \tag{7}$$

where M^* is the effective nucleon mass. From the equation of motion for the scalar meson, M^* is given by

$$M^* = M - \Phi = M - \frac{C_s^2}{M^2}(\rho_s - BM\Phi^2 - C\Phi^3), \tag{8}$$

where $C_s = g_s M/m_s$, and Φ/g_s is the ground-state expectation value of the field ϕ. The pressure P and energy density ϵ are also given as

$$P = P' + \frac{C_v^2}{2M^2}\rho^2 - U(\Phi) \tag{9}$$

and

$$\epsilon = \frac{\epsilon'}{1+V_n\rho'} + \frac{C_v^2}{2M^2}\rho^2 + U(\Phi) \tag{10}$$

, respectively, where $C_v = g_v M/m_v$,

$$P'(k_F, M^*) = \frac{\lambda}{12\pi^2}\{E_F^* k_F(E_F^{*2} - \frac{5}{2}M^{*2}) + \frac{3}{2}M^{*4}\ln(\frac{E_F^* + k_F}{M^*})\}, \tag{11}$$

and

$$\epsilon'(k_F, M^*) = E_F^*\rho' - P'. \tag{12}$$

with $E_F^* = \sqrt{k_F^2 + M^{*2}}$. We remark that the pressure and the energy density of free point-like nucleons can be given by eqs. (11) and (12), respectively, if we replace M^* by the free nucleon mass M. The baryonic chemical potential is also given by

$$\mu = E_F^* + V_n P'(k_F, M^*) + \frac{C_V^2}{M^2}\rho. \tag{13}$$

It is easy to check that the equations (9),(10) and (13) satisfy the thermodynamical identity [18], $\mu = (\epsilon + P)/\rho$, which yields the following relation,

$$M_0^* = \sqrt{E_{F0}^{*2} - k_{F0}^2} = [\{M - a_1 - V_n P'(k_{F0}, M_0^*) - C_v^2\rho_0/M^2\}^2 - k_{F0}^2]^{1/2} \tag{14}$$

at the normal density ρ_0, where a_1 is the binding energy of the normal nuclear matter and the subscripts "0" denotes "at the normal density". Since $P'(k_{F0}, M_0^*) > 0$, $\rho_0 > 0$ and $C_v^2 > 0$, it is shown that

$$M_0^* < \{(M - a_1)^2 - k_{F0}^2\}^{1/2}. \tag{15}$$

This condition gives the upper bound for M_0^*. For example, if we put $k_{F0} = 1.4 fm^{-1}$, we get $M_0^* < 0.94M$. (We remark that, due to the modification of baryon density (see eq. (5)), k_{F0} is larger in the theory with the EVE than in the theory of the point-like nucleons, for the given ρ_0.)

The incompressibility K at the normal density is defined as

$$K = 9\rho_0^2\frac{\partial^2 e}{\partial\rho^2}|_{\rho=\rho_0} = 9\frac{\partial P}{\partial\rho}|_{\rho=\rho_0} = 9\rho_0\frac{\partial\mu}{\partial\rho}|_{\rho=\rho_0}, \tag{16}$$

where $e = \epsilon/\rho$. In this paper, we also calculate the skewness coefficient, i.e., the third-order derivative K' of nuclear saturation curve. In our definition, K' is defined as

$$K' = 3\rho_0^3\frac{d^3 e}{d\rho^3}|_{\rho=\rho_0} = 3\rho_0\frac{d^2 P}{d\rho^2}|_{\rho=\rho_0} - \frac{4}{3}K = 3\rho_0^2\frac{d^2\mu}{d\rho^2}|_{\rho=\rho_0} - K. \tag{17}$$

In this definition, large K' means that the equations of state (EOS) is stiff at high density.

There is the eight independent parameters in this model: i.e., M, ρ_0, a_1, $R_n(=(3V_n/4\pi)^{1/3})$, C_s, C_v, B and C. In our calculations, we put $M=939$MeV, $\rho_0=0.17\text{fm}^{-3}$ and $a_1=16$MeV. The other five parameters R_n, C_s, C_v, B and C are determined phenomenologically. Besides the two conditions for the saturation (i.e., $e_0=M-a_1$ and $P=0$), if M_0^*, K and K' are given, we can determine the five parameters of the model. Conversely, we can calculate M_0^*, K and K', if the five parameters are given. As is seen in eq. (14), M_0^* is determined if R_n and C_v are given. Therefore, if we give one (two) quantity (quantities) in R_n, C_v and M_0^* and give two (one) quantities (quantity) in C_s, B, C, K and K', the other five quantities are automatically determined. We also remark that we put $R_n=0\sim 0.9$fm in our calculations. It is difficult to do the calculations at very large $R_n(>0.9\text{fm})$, which is close to $R_0(=\{3/(4\pi\rho_0)\}^{1/3}\sim 1.1\text{fm})$, i.e., the half of the averaged distance between two nucleons in the normal nuclear matter.

First we give R_n, M_0^* and K, and calculate K'. In fig. 1, we show K' as a function of R_n for the fixed values of K and M_0^*. In the cases of small $K(\leq 200\text{MeV})$, K' decreases in the large $R_n(\gtrsim 0.75\text{fm})$ region. We remark that, in the case of $K=200MeV$ and $M_0^*=0.9M$, K' does not decrease much as in the other cases in fig.1(a). K' also decreases in the large R_n region, in the cases of large $K(\geq 300\text{MeV})$ and small $M_0^*(=0.5M)$, although the absolute value of the decrease is not large as in the cases of small K. In the case of large K and the large $M_0^*(=0.9M)$, K' hardly decreases. The decrease of K' is related to the fact that the coefficient C, which has an important role for determining the high density behavior of EOS, becomes (more) negative in the large R_n region. C becomes (more) negative (i.e., attractive) to cancel the repulsive effect of the EVE, to realize the fixed value of K, as R_n increases. (The decrease of C is conspicuous in the cases of the small M_0^*, because C must also cancel the repulsive effects of the small M_0^* as well as the effects of the EVE. Also in the cases of small K, the decrease of C is conspicuous, because the small K must be reproduced.) As a result, K' becomes smaller in the large R_n region, because of the much negative C. It is also seen that K' always increases in the region of $R_n=0\sim 0.7$fm, as R_n increases, for the large $K(\geq 300\text{MeV})$. In these cases, the absolute value of the increase is large for the large $M_0^*(=0.9M)$. Therefore, K' is larger at $R_n=0.8$fm, which is often used as the nucleon radius, than at $R_n=0$, in the cases of the large K and the large M_0^*. On the contrary, K' is smaller at $R_n=0.8$fm than at $R_n=0$, in the cases of the small K and the small M_0^*.

Fig. 1(a),(b)

After K' is determined, we can also calculate Coulomb coefficient K_c in the leptodermous expansion (1), using the scaling model, i.e., using the following

equation [1],

$$K_c = -\frac{3q_{el}^2}{5R_0}\Big(\frac{9K'}{K} + 8\Big), \tag{18}$$

where q_{el} is the electric charge of proton. It is easily seen that K_c becomes more negative as K' increases in eq. (18), if K is fixed. K_c is negative and the absolute value is large in the EOS, which becomes stiffer at high density. In fig. 2, we show K_c as a function of K with the fixed values of R_n and M_0^*. In the case of $M_0^*/M = 0.5(0.7, 0.9)$, K_c is smaller (more negative) at $R_n = 0.8$fm than at $R_n = 0$, if $K \gtrsim 250(260, 160)$MeV. The reason is that, in these cases, K' is larger at $R_n = 0.8$fm than at $R_n = 0$, as is seen in fig. 1. It is remarkable that, in the case of $M_0^*/M = 0.5(0.7, 0.9)$, the K_c is larger (less negative) at $R_n = 0.8$fm than at $R_n = 0$, if $K \lesssim 250(260, 160)$MeV. Naturally, the reason is that, in these cases, K' is smaller at $R_n = 0.8$fm than at $R_n = 0$, as is also seen in fig. 1. The EVE make K' smaller and make K_c less negative for the small K. This change of K_c makes it possible to fit the empirical values of K_c for the small K. We also remark that the change of K_c by the EVE is opposite to the one by the vector meson self-interaction (VSI) [12]. It is reasonable that the repulsive EVE has opposite effects to the ones by the attractive VSI.

Fig. 2

If we give K, K_c and R_n, the other parameters of the model are also determined. (We remark that K' is also uniquely determined by eq. (18), if K and K_c are given.) In table 2, we show the examples of the parameters sets which reproduce the empirical values of K and K_c in table 1 with several values of R_n. For simplicity, we use the average values of K and K_c in table 1.

Table 2(a),(b),(c),(d)

We could not find the parameter sets which reproduce $(K, K_c) = (200.0, 2.577)$MeV at any R_n, as in the case of $R_n = 0$ [11], where the calculated K_c is always smaller than the empirical value 2.577MeV. This fact is understood as follows. At $R_n = 0$ and $K = 200$MeV, the largest M_0^* has the largest K_c which is closest to the empirical value. However, as is seen in fig. 1(a), K' does not become much smaller in the cases of $K = 200$MeV and large $M_0^*(\gtrsim 0.9M)$, even if we use a very large $R_n(> 0.8)$fm. As a result, in those cases, K_c hardly becomes large and does not reproduce the empirical value. On the other hand, in the cases of $K = 200$MeV and $M_0^* \lesssim 0.7M$, the value of K_c is much smaller at $R_n = 0$ than the empirical value. In those cases, although the EVE makes K' much smaller and makes K_c much larger in the large R_n region, K_c is still far from 2.577MeV.

In the case of $(K, K_c) = (250.0, -0.7065)$MeV, the empirical values is reproduced by this model in the regions of $R_n \approx 0 \sim 0.6$fm and $R_n = 0.85 \sim 0.88$fm.

The calculated effective nucleon mass M_0^* increases as R_n increases in the former region. The solution of parameter set disappears at $R_n = 0.6$fm, because of the upper bound condition (15) for M_0^*. The solution appears again at $R_n = 0.85$fm. This reappearance of the solutions is related to the fact that K' decreases in the large R_n region. The disappearance of the solution in the intermediate region of R_n occurs also in the cases of (K, K_c) =(300.0, -3.990)MeV and (350.0,-7.274)MeV. The solution disappears at $R_n = 0.72(0.79)$fm and reappears at $R_n = 0.86(0.85)$fm in the case of (K, K_c) =(300.0, -3.990)MeV ((350.0,-7.274)MeV).

In the cases of (K, K_c)=(150.0,5.861)MeV and (143,3.04)MeV, at $R_n = 0$, there is no parameter set, which reproduce the empirical values. However, we could find the parameter sets for (K, K_c) =(150.0,5.861)MeV and for(K, K_c)=(143, 3.04)MeV, if we put $R_n = 0.81 \sim 0.84$fm and $R_n = 0.76 \sim 0.87$fm, respectively. (See EOS10 and EOS11 in table 2(d).) The reason is that, as is seen in fig. 1, the EVE makes K' much smaller in the large R_n($\lesssim$0.75fm) region, in the case of $K \sim 150$MeV.

We remark that the coefficient C is always negative in the these solutions, as is seen in the EOS 10 and EOS 11. C becomes (more) negative (i.e., attractive) to cancel the large repulsive effects of the large R_n, as is mentioned before. Negative C may cause the difficulty such as a bifurcation of solution of Φ. However, this difficulty is modified by introducing the higher terms of Φ which hardly affect the nuclear matter properties at the normal density. For example, we add the following terms to $U(\Phi)$.

$$U_{56}(\Phi) = \frac{D}{5M}\Phi^5 + \frac{E}{6M^2}\Phi^6, \tag{19}$$

where D and E are the dimensionless constants. If we put $D = -0.00755$ and $E = 0.006$, for example, we get the EOS 12 in table 2(d). Although, at $\rho = \rho_0$, the EOS 12 has almost the same properties as in the EOS 11, it has only one solution as is seen fig 3, where $dU(\Phi)/d\Phi - \rho_s$ is shown as a function of Φ, both in the cases of the EOS 11 and of the EOS 12. (We remark $dU(\Phi)/d\Phi - \rho_s = 0$ is a equivalent condition to eq. (8).)

We also calculate the volume-symmetry coefficient K_{vs} in the expansion (1). Because the ρ-meson effects are important in the symmetry properties [21][22], we add the standard ρ-meson terms to the Lagrangian [21][22][19]. According to the modification, the following term is added to the energy density [21][22][19].

$$\epsilon_\rho = \frac{g_\rho^2}{8m_\rho^2}(\rho_p - \rho_n)^2 = \frac{C_\rho^2}{8M^2}\rho_3^2, \qquad \rho_3 = \rho_p - \rho_n, \tag{20}$$

where m_ρ, g_ρ, ρ_p and ρ_n are ρ meson mass, ρ-nucleon coupling, proton density and neutron density, respectively, and $C_\rho = g_\rho M/m_\rho$. In the theory with the

EVE [19], ρ_p and ρ_n are given by

$$\rho_p = \frac{\rho'_p}{1+V_n\rho'} \qquad \text{and} \qquad \rho_n = \frac{\rho'_n}{1+V_n\rho'}, \tag{21}$$

, where the expressions for ρ'_p and ρ'_n have the same expressions as the densities of the point-like proton and neutron, respectively. In the mean field theory, the inclusion of the ρ-meson effects do not affect the saturation conditions and do not change the properties such as K and K_c in the symmetric nuclear matter. Therefore, the determination of the parameters R_n, C_s, C_v, B and C in fitting the data of K and K_c is not affected by this modification of the Lagrangian. Using the modified Lagrangian with the parameter sets in table 2, we calculate K_{vs} with aid of the scaling model [1]: i.e.,

$$K_{vs} = K_{sym} - L\Big(9\frac{K'}{K} + 6\Big), \tag{22}$$

where

$$L = 3\rho_0\frac{da_4}{d\rho}|_{\rho=\rho_0}, \qquad K_{sym} = 9\rho_0^2\frac{d^2a_4}{d\rho^2}|_{\rho=\rho_0}, \qquad \text{and} \qquad a_4 = \frac{1}{2}\rho\frac{\partial^2\epsilon}{\partial\rho_3^2}|_{\rho_3=0}. \tag{23}$$

The results are also summarized in table 2. In these calculations, we determine the ρ meson coupling g_ρ so as to realize $a_4 = 30.0$MeV at $\rho = \rho_0$. By comparing the table 1 and table 2, it is seen that $K = 300 \pm 50$MeV is favorable to account for the empirical values of K, K_c and K_{vs} simultaneously, as in the case of $R_n = 0$ [11], and as in the case of the mean-field theory with the VSI [12]. This unchanged conclusion is related to the fact that K_{vs} is very sensitive to the ratio K'/K, which is adjusted to the empirical value in table 2. It seems that this feature is not changed drastically, if we use the relativistic mean-field theory and the scaling model.

In the last, we examine whether the QMC result is well reproduced by the EVE model. According to ref. [17], $K = 200$MeV and $M_0^* = 0.906M$, if the bag radius $R_B = 0.8$fm is used. Since the value of K' and K_c is not shown in the reference, we calculate K' by using the fig. 2 in ref. [17], where the saturation curve of the nuclear matter in the QMC model is shown. The result is $K' \sim -85$MeV. (This value corresponds to $K_c \sim -3.2$MeV, which is somewhat larger than the empirical value, 2.577 ± 2.06MeV.) We search the parameter set of the EVE model, which reproduces the $K = 200$MeV, $M_0^* = 0.906M$ and $K' = -85$MeV. The results are shown in table 3. In the parameter set, $R_n (= 0.568\text{fm})$ is somewhat smaller than $R_B (= 0.8\text{fm})$. The physical meaning of R_n in the EVE model may be somewhat different from that of the bag radius R_B in the QMC model.

Table 3

In summary, we have studied the compressional properties of nuclear matter by using the relativistic mean field theory with the nonlinear σ terms and the EVE of the nucleons, under the assumption of the scaling model. We found that the EVE yields the possibility to reproduce the empirical values of K_c with the small $K \sim 150$MeV. However, if we require that K_{vs} should be reproduced as well as K_c, $K = 300 \pm 50$MeV is favorable. It seem to be difficult to change this conclusion drastically, in the framework of the relativistic mean field theory and the scaling model.

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Table and Figure Captions

Table 1
The sets of the empirical values of K, K_c and K_{vs}. (Shown in MeV.) The sets $1 \sim 5$ is the data from the table 3 in ref. [2]. (According to the conclusion in ref. [2], we only show the data in the cases of $K = 150 \sim 350$MeV.) The set 6 is one example of the data of the table IV in ref. [3].

Table 2
Parameter sets fitted for the mean values of the empirical values of K and K_c in table 1. The K, K_c, K', K_{vs}, L and K_{sym} are shown in MeV, while R_n are shown in fm. In EOS12, D=-0.00755 and E=0.006. (See eq. (19).)

Table 3
Parameter sets fitted for $K = 200$MeV, $M_0^* = 0.906M$, and $K' = -85$MeV. R_n are shown in fm.

Fig. 1 K' as a function of R_n with several values of M_0^* and K. In the figure (a) ((b)), the solid line, the dotted line, and the dash-dotted line are the results with M_0^*/M =0.5, 0.7, and 0.9, respectively, for $K = 150(300)$MeV. In the figure (a) ((b)), the bold solid line, the bold dotted line, and the bold dash-dotted line are the results with M_0^*/M =0.5, 0.7, and 0.9, respectively, for $K = 200(400)$MeV.

Fig. 2 K_c as a function of K with fixed values of R_n and M_0^*. The solid line, the dotted line, and the dash-dotted line are the results with M_0^*/M =0.5, 0.7, and 0.9, respectively, in the case of $R_n = 0$. The bold solid line, the bold dotted line, and the bold dash-dotted line are the results with M_0^*/M =0.5, 0.7, and 0.9, respectively, in the case of $R_n = 0.8$fm. The crosses with error bars are the data sets $1 \sim 5$ in table 1. The solid squares are the data from the table IV in ref. [3]. (For simplicity of the figure, we omit the error bars in the latter data.)

Fig. 3 $dU(\Phi)/d\Phi - \rho_s$ as a function of Φ. The solid line and the dotted line are the results of EOS11 and EOS12, respectively.

	Set 1	Set 2	Set 3	Set 4	Set 5	Set 6
K	150.0	200.0	250.0	300.0	350.0	143 ± 53
K_c	5.861 ± 2.06	2.577 ± 2.06	-0.7065 ± 2.06	-3.990 ± 2.06	-7.274 ± 2.06	3.04 ± 4
K_{vs}	66.83 ± 101	-46.94 ± 101	-160.7 ± 101	-274.5 ± 101	-388.3 ± 101	34 ± 159

Table 1

R_n	C_s^2	C_v^2	B	C
0.568	106.17	16.472	-2.589$\times 10^{-2}$	1.049

Table 3

EOS	1	2	3
K	250.0	300.0	350.0
K_c	-0.7065	-3.990	-7.274
K'	-197	-94.0	56.0
R_n	0.00	0.00	0.00
M_0^*/M	0.910	0.831	0.609
K_{vs}	53.83	-275.5	-626.8
L	76.61	78.90	93.81
K_{sym}	-29.20	-24.68	71.27
C_s^2	42.392	144.56	293.29
C_v^2	18.459	66.337	197.85
B	-0.5282	-5.067$\times10^{-3}$	1.595$\times10^{-3}$
C	5.071	0.1028	-1.686$\times10^{-3}$
C_ρ^2	90.31	84.11	59.94

Table 2(a)

EOS	4	5	6
K	250.0	300.0	350.0
K_c	-0.7065	-3.990	-7.274
K'	-197	-94.0	56.0
R_n	0.60	0.72	0.79
M_0^*/M	0.931	0.919	0.902
K_{vs}	77.41	-244.9	-600.1
L	79.75	84.81	91.54
K_{sym}	-9.193	24.76	80.95
C_s^2	18.121	55.348	100.08
C_v^2	0.42953	2.0916	6.7682
B	-2.063	-0.2924	3.295$\times10^{-2}$
C	23.00	3.728	0.2168
C_ρ^2	83.83	75.39	66.67

Table 2(b)

EOS	7	8	9
K	250.0	300.0	350.0
K_c	-0.7065	-3.990	-7.274
K'	-197	-94.0	56.0
R_n	0.85	0.86	0.85
M_0^*/M	0.573	0.540	0.643
K_{vs}	1054	597.5	-407.3
L	150.5	164.1	133.1
K_{sym}	890.6	1119	582.8
C_s^2	289.99	302.02	245.09
C_v^2	181.40	195.55	144.14
B	2.785×10^{-3}	2.213×10^{-3}	3.421×10^{-3}
C	-7.121×10^{-3}	-5.987×10^{-3}	-9.667×10^{-3}
C_ρ^2	11.18	2.037	23.87

Table 2(c)

EOS	10	11	12
K	150.0	143.0	143.0
K_c	5.861	3.04	3.04
K'	-260	-190	-190
R_n	0.81	0.76	0.76
M_0^*/M	0.574	0.623	0.623
K_{vs}	1988	1060	1060
L	137.4	116.9	116.9
K_{sym}	668.8	364.5	364.5
C_s^2	307.94	291.23	287.83
C_v^2	189.47	170.39	170.39
B	3.370×10^{-3}	4.527×10^{-3}	3.881×10^{-3}
C	-7.261×10^{-3}	-8.976×10^{-3}	-5.556×10^{-3}
C_ρ^2	19.83	36.22	36.22

Table 2(d)